

Inert Scalar Doublet Asymmetry as Origin of Dark Matter

Mikaël Dhen and Thomas Hambye^{1,*}

¹*Service de Physique Théorique
Université Libre de Bruxelles
Boulevard du Triomphe, CP225, 1050 Brussels, Belgium*

In the inert scalar doublet framework, we analyze what would be the effect of a $B-L$ asymmetry that could have been produced in the Universe thermal bath at high temperature. We show that, unless the “ λ_5 ” scalar interaction is tiny, this asymmetry is automatically reprocessed in part into an inert scalar asymmetry that could be at the origin of dark matter today. Along this scenario, the inert mass scale lies in the few-TeV range and direct detection constraints require that the inert scalar particles decay into a lighter dark matter particle which, as the inert doublet, is odd under a Z_2 symmetry.

I. INTRODUCTION

The similarity of baryonic and dark matter (DM) abundances, determined by the observation of the Cosmic Microwave Background (CMB) anisotropies, $\Omega_{DM}/\Omega_B = 5.4 \pm 0.1$ [1], has motivated a long series of scenarios where both abundances have a related or even very same origin. Since the baryon asymmetry is to a very good approximation totally asymmetric – no primordial population of antibaryons has been observed in the Universe – a common origin of both abundances suggests that the DM abundance today would be associated to the generation of a DM particle-antiparticle asymmetry (see e.g. the reviews of Refs. [2–5]). In the following, we will show how this can be realized from the generation of a scalar inert doublet asymmetry. The inert scalar doublet DM framework (IDM) [6–9] simply consists in adding to the Standard Model (SM) a single scalar doublet, H_2 , odd under a Z_2 symmetry. The most general scalar potential is in this case

$$V = m_1^2 |H_1|^2 + m_2^2 |H_2|^2 + \lambda_1 |H_1|^4 + \lambda_2 |H_2|^4 + \lambda_3 |H_1|^2 |H_2|^2 + \lambda_4 |H_1^\dagger H_2|^2 + \frac{\lambda_5}{2} \left[(H_1^\dagger H_2)^2 + h.c. \right], \quad (1)$$

where the SM and the inert scalar doublets can be written as

$$H_1 = \begin{pmatrix} \phi^+ \\ v/\sqrt{2} + \phi^0 \end{pmatrix} \quad \text{and} \quad H_2 = \begin{pmatrix} \eta^+ \\ \eta^0 \end{pmatrix}, \quad (2)$$

with $\phi^0 \equiv (h + i\phi^3)/\sqrt{2}$ and $\eta^0 \equiv (H^0 + iA^0)/\sqrt{2}$. In the scalar potential, m_2^2 is assumed to be positive to insure that H_2 doesn’t acquire a vev, so that its lightest (neutral) component is stable, unless there exists a lighter Z_2 odd particle into which it can decay. This is the possibility we will ultimately consider in the following, in order to satisfy the direct detection constraints, see Section IV. But before, up to the end of Section III, let’s restrict ourselves to the usual IDM setup and see what happens in this minimal framework.

Prior to electroweak symmetry breaking (EWSB), all H_2 components have mass $m_{H_2} = m_2$, whereas after EWSB ($v =$

246 GeV), they get split in mass

$$m_{H^0}^2 = m_2^2 + \lambda_{H^0} v^2, \quad m_{A^0}^2 = m_2^2 + \lambda_{A^0} v^2, \quad m_{\eta^\pm}^2 = m_2^2 + \lambda_{H^\pm} v^2, \quad (3)$$

with $\lambda_{H^\pm} = \lambda_3/2$ and $\lambda_{H^0, A^0} = (\lambda_3 + \lambda_4 \pm \lambda_5)/2$. In the following, we will assume, without loss of generality, that λ_5 is negative, so that H^0 is lighter than A^0 . Various well-known constraints hold on the parameters of the theory. Tree level vacuum stability requires $\lambda_{1,2} > 0$, $\lambda_{H^0, A^0, H^\pm} > -\sqrt{\lambda_1 \lambda_2} \approx -0.36\sqrt{\lambda_2}$. EW precision test observables require $\Delta T \simeq (m_{\eta^\pm} - m_{A^0})(m_{\eta^\pm} - m_{H^0})/12\pi^2\alpha v^2 \lesssim 10^{-1}$. Z decay width constraint at LEP requires $m_{A^0} + m_{H^0} > m_Z$ and $m_{\eta^\pm} > m_Z/2$. Direct detection constraint importantly requires that the Z exchange diagram is kinematically forbidden, i.e. $m_{A^0} - m_{H^0} \gtrsim \mu_r \beta_{DM}^2/2$, where β_{DM} is the DM halo velocity with respect to the earth, and $\mu_r = m_{H^0} m_{\mathcal{N}}/(m_{H^0} + m_{\mathcal{N}})$ is the reduced mass of the system for the nucleus \mathcal{N} used by the experiment. For $m_{H^0} \gg m_{\mathcal{N}}$ and Xenon nucleus, using an average velocity of ~ 270 km/s, this constraint can be rephrased as $m_{A^0} - m_{H^0} \gtrsim \delta m_{min} \sim 50$ keV. Taking into account the velocity distribution around this central value, and the recoil energy sensitivity of the experiments, the minimum splitting becomes $\delta m_{min} \sim 180$ keV, although a more robust constraint is $\delta m_{min} \sim 100$ keV [10], which translates as

$$|\lambda_5| \gtrsim 3.3 \cdot 10^{-6} \cdot \left(\frac{m_{H^0}}{\text{TeV}} \right) \cdot \left(\frac{\delta m_{min}}{100 \text{ keV}} \right). \quad (4)$$

It is well known that the IDM can account for the observed DM relic abundance via the usual freeze-out mechanism, and be in agreement with direct detection constraints, for DM masses in the ranges $\sim [50, 80]$ GeV and above ~ 540 GeV, up to the ~ 40 - 50 TeV unitarity bound [6, 7, 9, 11, 12]. As we will show, it could also be responsible for the DM relic density in an asymmetric way.

II. ASYMMETRIC PRODUCTION OF THE DM RELIC DENSITY

As said above, in this Section and the following one we stick to the minimal IDM framework, as defined in the introduction. Let us make two simple starting assumptions. First,

*Electronic address: mikadhen@ulb.ac.be; thambye@ulb.ac.be

let us assume that the symmetric component of the relic density left after freeze-out is smaller than the observed value. Fast SM gauge scatterings automatically care for that for m_{H^0} within the $\sim 120 - 540$ GeV range, whereas for other values of m_{H^0} large enough $\lambda_{3,4}$ interactions can take care of that [9]. This implies a symmetric annihilation cross section larger than the usual thermal freeze-out value ~ 1 pb, which means a freeze-out temperature T_{fo} smaller than the usual $T_{fo} \sim m_{H^0}/25$ value. Second, let us assume that a $B-L$ asymmetry has been generated at a temperature T_{B-L} above m_{H_2} and above the EWPT temperature T_{EW} (which we take as the temperature where the vacuum expectation value of the SM scalar field becomes sizable, that is $T_{EW} \approx 165$ GeV from Ref. [13]). We do not care about the way this $B-L$ asymmetry could have been generated. It could be due for example to the straightforward leptogenesis mechanism. Note that, as well-known, if a $B-L$ asymmetry is generated at high temperature, a $H_1-H_1^*$ asymmetry will also be created automatically at high temperature from thermal equilibrium SM interactions [14].

If a $B-L$ (and thus H_1 asymmetry) is created at high temperature, an inert doublet $H_2-H_2^*$ asymmetry is to be expected too. The scalar potential of Eq. (1) contains the λ_5 interaction which uniquely does not conserve the number of H_2 minus the number of H_2^* (as well as the number of H_1 minus the number of H_1^*). This interaction is in thermal equilibrium at $T \sim m_{H_2}$ if at this temperature the corresponding Γ_{λ_5} scattering rate, given in the Appendix, is larger than the Hubble rate, which gives the condition

$$|\lambda_5| \gtrsim 10^{-6} \cdot (m_{H_2}/\text{TeV})^{1/2}. \quad (5)$$

If Eq. (5) is satisfied, the λ_5 interaction equilibrates the H_2 and H_1 (and $B-L$) asymmetries.¹ In particular even if, as we assume here, no H_2 asymmetry is created at high energies, such an asymmetry will be created anyway as soon as the $B-L$ asymmetry is created. In other words, the inert DM model contains an interaction which basically implies that “Higgsogenesis” [16] production of a DM asymmetry is at work.² Note that the scenario could work also the other way around, i.e. a primordial DM asymmetry could be at the origin of baryogenesis via the same λ_5 equilibration interaction, a possibility we will not consider here (for a scenario of this kind see Ref. [19]).

¹ Actually, in the few-TeV asymmetric inert DM scenario considered in Ref. [15], it is assumed instead that the λ_5 interaction could have never been in thermal equilibrium. In this case, the DM asymmetry would have been created explicitly at high energies, basically independently of the $B-L$ asymmetry. If the inert scalar is the DM particle, it turns out that the lower bound of Eq. (4) implies that Eq. (5) must anyway hold for TeV masses. For instance, Eq. (4) with a 100 keV (180 keV) mass splitting implies Eq. (5) for $m_{H_2} \gtrsim 100$ GeV (30 GeV).

² In Ref. [16], a X_1 fermion singlet DM framework is considered with an extra X_2 fermion doublet and an X -symmetry. An X_2 asymmetry is created from a X -symmetry violating $X_2^2 H_1^2$ non-renormalizable interaction, which is afterwards reprocessed into a X_1 symmetry through X_2 decays. Asymmetric frameworks based on the equilibration of the SM scalar asymmetry with a dark sector asymmetry, based on several new dark sector particles, or based on various possibilities of a $SU(2)_L$ multiplet, can also be found in Refs. [17] and [18] respectively.

In the following, we will consider in details and chronologically what happens when the temperature of the Universe cools down from $T \gg m_{H_2}$ to today $T \ll T_{EW}$, crossing $m_{H_2} > T_{\lambda_5} > T_{fo} > T_{EW}$, with T_{λ_5} the temperature where the scattering induced by the λ_5 interaction decouples and T_{fo} the freeze-out temperature at which the total annihilation cross section decouples. Given that the inert doublet components undergo gauge interactions, T_{λ_5} is sizably larger than T_{fo} , unless λ_5 is of order one, which as we will see is not a viable option for the case we are interested in (where the DM asymmetry is responsible for most of the relic density). Similarly, as we will see, $T_{fo} \gtrsim T_{EW}$, i.e. few-TeV DM, is also generically necessary in order to have a viable scenario (as in the scenario of Ref. [15]), but some violation of this inequality is possible. A sketch of the scenario, applied to our framework, is shown in Fig. 1.

A. $T \gtrsim m_{H_2}$

At temperature above T_{EW} , all 4 inert doublet components have a common mass $m_{H_2} = m_2$. If Eq. (5) is satisfied, the chemical potential of both scalar doublets are equal, $\mu_{H_2} = \mu_{H_1}$. Together with the usual SM chemical equilibrium relations (from thermal equilibrium SM processes [14]), the $\mu_{\eta^+} = \mu_{\eta^0}$ relation (from e.g. $\eta^+ \eta^{0*} \leftrightarrow SM$ processes), and the hypercharge relation

$$\sum_i (\Delta_{Q_i} + 4\Delta_{u_i} - 2\Delta_{d_i} - \Delta_{\ell_i} - 2\Delta_{e_i}) + \Delta_{H_1} + \Delta_{H_2} = 0, \quad (6)$$

it simply gives

$$\Delta_{H_2} = \Delta_{H_1} = -\frac{4}{23} \Delta_{B-L}. \quad (7)$$

From now on, we define for each species X the asymmetry $\Delta_X \equiv Y_X - Y_{\bar{X}}$ and the total density $\Sigma_X \equiv Y_X + Y_{\bar{X}}$, where $Y_X \equiv n_X/s$ is the particle number density-to-entropy ratio of X . Since we are dealing with asymmetries, we also define the number of degrees of freedom by summing the number of particles (or antiparticles but not both), i.e. $g_X = 1$ for a $SU(2)_L$ singlet, and $g_X = 2$ for a doublet.

As well-known, for similar $B-L$ and DM asymmetries, the DM relic density constraint requires m_{DM} to have a mass of few GeV (more exactly, from Eq. (7) and taking into account the Y_{B-L} to Y_B ratio which holds in this case, Eq. (21) below, one would need $m_{DM} \approx 10$ GeV). As this possibility is excluded by collider constraints, this implies that a subsequent suppression of the DM asymmetry by a factor of $\sim (10 \text{ GeV}/m_{DM})$ must necessarily occur. Two different types of suppressions can naturally take place. A first one is a Boltzmann suppression from asymmetry violating scatterings, used in several other DM models, see e.g. Refs. [16, 20]. In our scenario, it can arise from the λ_5 interaction within the period $m_{H_2} > T > T_{\lambda_5}$. The other possible suppression can arise later when $T \lesssim T_{EW}$ from the combined effect of DM oscillations and symmetric annihilations.

B. $m_{H_2} \gtrsim T \geq T_{\lambda_5}$

Once the temperature drops below m_{H_2} , if the λ_5 interaction goes on to be in thermal equilibrium, the H_2 asymmetry gets Boltzmann suppressed. This can be directly seen from the Boltzmann equation of Δ_{H_2} , valid for $T \geq T_{EW}$,

$$\frac{d\Delta_{H_2}}{dz} = -\frac{4}{sHz} \left(\frac{\Delta_{H_2}}{Y_{H_2}^{eq}} - \frac{\Delta_{H_1}}{Y_{H_1}^{eq}} \right) \gamma_{\lambda_5}, \quad (8)$$

where $z \equiv m_{H_2}/T$, $H(z)$ is the Hubble rate and $\gamma_{\lambda_5}(z)$ is the reaction density of the λ_5 scatterings, given in Appendix. It includes both pair annihilation/creation $H_2 H_2 \leftrightarrow H_1 H_1$ and “t-channel” $H_2 \bar{H}_1 \leftrightarrow H_1 \bar{H}_2$ processes. The λ_5 interaction leaves intact the sum of the asymmetries of H_1 and H_2 but not each asymmetry individually. Once T drops below m_{H_2} , the first term in the r.h.s. of Eq. (8) is enhanced with respect to the second term by the fact that $Y_{H_2}^{eq}$ is Boltzmann suppressed, unlike $Y_{H_1}^{eq}$. This Boltzmann suppression of the asymmetry lasts until the λ_5 induced scatterings decouple, at $T = T_{\lambda_5}$, when $\Gamma_{\lambda_5} \simeq H$. Quantitatively, this can be accounted by the usual k -factor which gives the asymmetry as a function of the temperature

$$\Delta_{H_2} = \frac{T^2}{6s} g_{H_2} \mu_{H_2} k(z), \quad (9)$$

$$k(z) \equiv \frac{6}{4\pi^2} \int_0^\infty x^2 \sinh^{-2} \left(\frac{1}{2} \sqrt{x^2 + z^2} \right) dx. \quad (10)$$

Down to T_{λ_5} , the chemical potential relation $\mu_{H_2} = \mu_{H_1}$ still holds, and we get

$$\Delta_{H_2}(z) = \frac{k(z)}{2} \Delta_{H_1}. \quad (11)$$

This is nothing but the solution which makes the r.h.s. of Eq. (8) to vanish. Together with the other chemical relations above, the asymmetry reads at $T = T_{\lambda_5}$ (similarly to the fermion doublet case of Ref. [16])

$$\Delta_{H_2}(z_{\lambda_5}) = -\frac{16k(z_{\lambda_5})}{158 + 13k(z_{\lambda_5})} \Delta_{B-L}. \quad (12)$$

For practical reasons, it is convenient to define

$$\Delta_{H_2}^{\lambda_5} \equiv |\Delta_{H_2}(z_{\lambda_5})|. \quad (13)$$

Let us note that since $z_{\lambda_5} \gg 1$, the k -factor can be approximated by

$$k(z_{\lambda_5}) \simeq 12 \left(\frac{z_{\lambda_5}}{2\pi} \right)^{3/2} e^{-z_{\lambda_5}}. \quad (14)$$

Clearly, the λ_5 coupling must not be too large in order to avoid a too strong exponential suppression of the H_2 asymmetry.

As emphasized in Ref. [16], for fermion quartic interactions, the last λ_5 induced channels to decouple are the t-channel ones,³ simply because they are less Boltzmann suppressed than the other ones. The value of z_{λ_5} is given by

the condition that the Γ_{λ_5} rate is equal to the Hubble rate H . One has for example that if $m_{H_1} = 10$ TeV, $z_{\lambda_5} = 2(15)$ for $\lambda_5 = 3 \cdot 10^{-6}(5 \cdot 10^{-6})$. For somewhat smaller values of λ_5 , even if the λ_5 reactions doesn't enter in thermal equilibrium (as defined by Eq. (5)), a numerical integration of the Boltzmann equations shows that still a number of scattering processes occur nevertheless, what can lead to a sizable asymmetry.

C. $T_{\lambda_5} > T > T_{f0}$

During this period, the Δ_{H_2} asymmetry stays constant unlike the total abundance Σ_{H_2} , whose Boltzmann equation for this period reads

$$\frac{d\Sigma_{H_2}}{dz} = -\frac{\langle \sigma_{eff} v \rangle s}{zH} \left[\Sigma_{H_2}^2 - \Delta_{H_2}^{\lambda_5 2} - \Sigma_{H_2}^{eq 2} \right], \quad (15)$$

where $\langle \sigma_{eff} v \rangle$ is the effective thermal cross section of the $H_2 \bar{H}_2 \leftrightarrow SM SM$ annihilations, given in the Appendix. With a constant $\Delta_{H_2}^{\lambda_5}$, as it is the case during this period, the solution of Eq. (15) at freeze-out is to a good approximation given by

$$\Sigma_{H_2}(z_{fo}) \simeq \left[\Delta_{H_2}^{\lambda_5 2} + \Sigma_{H_2}^{eq 2}(z_{fo}) \right]^{1/2}, \quad (16)$$

which is nothing but the expression which makes the r.h.s. of Eq. (15) to vanish. Here, by z_{fo} we mean the usual freeze-out value given by the equation

$$z_{fo} \simeq \ln \frac{0.0038 \cdot m_{Pl} 2g_{H_2} m_{H_2} \langle \sigma_{eff} v \rangle}{\sqrt{g_*} z_{fo}}. \quad (17)$$

If the annihilations are fast enough to leave at T_{fo} a symmetric component smaller than the asymmetric one (which is typically satisfied for $\langle \sigma_{eff} v \rangle \gtrsim 1$ pb), the following relation holds, $\Sigma_{H_2}(z_{fo}) \simeq \Delta_{H_2}^{\lambda_5} \gg \Sigma_{H_2}^{eq}(z_{fo})$. Given the sign of the baryon asymmetry, this means at T_{fo} , $\Sigma_{H_2} \sim -\Delta_{H_2} \sim -Y_{\bar{H}_2} \gg Y_{H_2}, \Sigma_{H_2}^{eq}$.

D. $T_{f0} > T > T_{EW}$

Nothing is expected to happen during this period. The H_2 total density left at T_{fo} is left intact until T_{EW} , temperature at which the total density and asymmetry are given by (for a dominant asymmetric component)

$$\Sigma_{H_2}(z_{EW}) \simeq \Sigma_{H_2}(z_{fo}) \simeq \Delta_{H_2}^{\lambda_5} \quad \text{and} \quad |\Delta_{H_2}(z_{EW})| = \Delta_{H_2}^{\lambda_5}. \quad (18)$$

E. $T < T_{EW}$

Next, once the temperature drops below T_{EW} , two new effects enter into play: generation of mass splittings between

³ We thank G. Servant and S. Tulin for discussions on the importance of these channels.

the H^0 , A^0 and η^+ components and possibly fast inert particle-antiparticle oscillations $\eta^0 \leftrightarrow \eta^{0*}$. The effect of the mass splittings generated by the SM scalar vev, Eq. (3), is of moderate importance. Assuming, as said above, that the H^0 component is the lightest one (i.e. $\lambda_5 < 0$), they imply that the other components will ultimately decay to H^0 . But these decays conserve the number of inert scalar particles. They just convert the H_2 asymmetry created before EWSB (with mass m_{H_2}) into a DM relic density of selfconjugated DM particles H^0 (with mass $m_{H^0} = m_{DM}$, different from m_{H_2} unless λ_{H^0} vanishes). More important is the potential effect of the much faster inert particle-antiparticle oscillations $\eta^0 \leftrightarrow \eta^{0*}$ caused by the λ_5 interactions. The rate of an oscillating particle is simply given by the value of the associated mass splitting [15, 21, 22], i.e. $\Gamma_{osc} = \delta m = m_{A^0} - m_{H^0}$. For $T < T_{EW}$, this rate is very fast compared to the Hubble rate

$$\frac{\Gamma_{osc}}{H(T)} \simeq 2 \cdot 10^{15} \cdot |\lambda_5| \cdot \left(\frac{100 \text{ GeV}}{T} \right)^2 \cdot \left(\frac{\text{TeV}}{m_{H^0}} \right). \quad (19)$$

The effect of these oscillations depends obviously on whether they do occur, which Eq. (19) doesn't necessarily imply, and on whether symmetric annihilations do occur after EWSB. Actually, even if the freeze-out occurs before EWSB, this does not imply that symmetric annihilations could not restart again after EWSB, due to oscillations [21, 22]. This could easily be the case because, even if one starts with a pure asymmetry, the oscillations will quickly give a number density of each population much larger than their thermal equilibrium values, roughly $n_{\eta^0} \sim n_{\eta^{0*}} \sim |\Delta n_{\eta^0}|/2 \gg n_{\eta^0}^{eq}$, so that $|\Delta n_{\eta^0}| \langle \sigma v \rangle > H$ can hold even if $n_{\eta^0}^{eq} \langle \sigma v \rangle < H$. If these annihilations occur, they will anyway reduce the DM abundance, as no inverse processes will occur in this case. Let us consider both possible cases separately.

1. $T < T_{EW}$: No symmetric annihilations after EWSB

If no symmetric annihilations arise after EWSB, oscillations have simply no effect. They quickly reconvert a pure η^0 population, or a pure η^{0*} population, into an oscillating mixed η^0 - η^{0*} population, but they do not change the number of inert states [21]. In this case, the number of H^0 particles left today will be simply equal to the number of inert scalar particles stored in the H_2 asymmetry before EWSB, i.e. $Y_{DM}^{today} \simeq \Delta_{H_2}^{\lambda_5}$, that means the H^0 density is equal to the asymmetry left after λ_5 interaction's decoupling. From Eqs. (12) and (18), this gives

$$Y_{DM}^{today} = \frac{16k(z_{\lambda_5})}{158 + 13k_{H_2}(z_{\lambda_5})} \Delta_{B-L}, \quad (20)$$

with $k(z_{\lambda_5})$ given by Eq. (14). Only the relation between the value of Y_B today and Δ_{B-L} changes after EWSB, as a result of the fact that below T_{EW} the conservation of electric charge holds rather than conservation of Y and T_3 . We get

$$Y_B^{today} \simeq \frac{12}{37} \Delta_{B-L}. \quad (21)$$

As a result, the DM density reads

$$Y_{DM}^{today} = \frac{148 \cdot k(z_{\lambda_5})}{474 + 39 \cdot k(z_{\lambda_5})} \cdot Y_B^{today}, \quad (22)$$

and the actual DM to baryon density ratio is given by

$$\frac{\Omega_{DM}}{\Omega_B} = \frac{Y_{DM}^{today}}{Y_B^{today}} \cdot \left(\frac{m_{H^0}}{1 \text{ GeV}} \right) = \frac{148 \cdot k(z_{\lambda_5})}{474 + 39 \cdot k(z_{\lambda_5})} \cdot \left(\frac{m_{H^0}}{1 \text{ GeV}} \right). \quad (23)$$

This is the final result if no symmetric annihilations occur after EWSB. However, it must be stressed that it is not mandatory to avoid symmetric annihilations after EWSB. On the contrary, if the λ_5 interaction above does not provide enough suppression, these scattering processes could easily provide it, without the need of any special tuning. This is what we will now quantify.

2. $T < T_{EW}$: Symmetric annihilations after EWSB

Possible effects of dark matter oscillations have been studied in Refs. [21, 22]. As said above, since oscillations reprocess the asymmetry into both particle and antiparticle densities, their main effect is to allow the symmetric annihilations to start again. Even if, as Eq. (19) shows, the oscillation rate is much faster than the Hubble rate, this doesn't necessarily mean that oscillations (and thus eventually annihilations) do occur. As shown in these references, if dark matter undergoes fast annihilations or elastic scatterings, these processes can break the coherence of the η^0 - η^{0*} states, preventing them from oscillating. The interplay of the oscillations with the other processes is actually more complicated in our doublet scenario than in the singlet setups considered in Refs. [21, 22]. On top of η^0 - η^{0*} annihilations and η^0 or η^{0*} elastic scatterings, there are charged η^\pm states, which at T_{EW} are responsible for half of the asymmetry and do not oscillate. Fast inelastic scatterings can change neutral states into these charged states and vice et versa. Moreover, as said above, all states ultimately become real H^0 states, which obviously do not oscillate.

Let us first consider what happens to the neutral states, as if there were no charged states. In this case, one has two important processes. On the one hand, there are $\eta^0 \eta^{0*} \rightarrow SSM$ annihilations processes which are dominated by their $\lambda_{3,4}$ interaction contribution. On the other hand there are $\eta^{0(*)} SM \rightarrow \eta^{0(*)} SM$ elastic scatterings which are dominated by their t-channel Z exchange contribution. The later dominates over the $\lambda_{3,4}$ elastic scattering contribution. For $T \gtrsim m_h$, this stems from the fact that it involves a t-channel mediator whose mass is much smaller than the inert scalar mass, $m_Z \ll m_{H_2}$. It scales as $\Gamma_{scat}^{gauge} \simeq G_F^2 T^5$ as compared to the quartic coupling elastic contribution which scales as $\Gamma_{scat}^{quartic} = n_{H_1}^{eq} \langle \sigma v \rangle \simeq \lambda_{3,4}^2 T^3 / m_{H_2}^2$. For $T < m_h$, the quartic contribution is also subleading because it is Boltzmann suppressed, unlike the gauge one. As a result, the gauge elastic contribution is the

last to decouple. The Z exchange process is relevant for preventing η^0 - η^{0*} states from oscillating because the gauge interaction is odd under η^0 - η^{0*} exchange [22]. Thus the relevant question is, down to which temperature will these processes effectively prevent the oscillations to start? At first sight, we could think that oscillations will start only once the scattering rate Γ_{scat} goes below the oscillation rate $\Gamma_{osc} = \delta m$ (if at this time both rates are still larger than the Hubble rate). This turns out to occur at a rather low temperature, $T_{osc} \sim \text{few GeV}$ scale. In this case one would be back to the “no symmetric annihilation” case above, because oscillations have practically no more effect at this temperature, where the annihilation rate is already largely suppressed. However, an integration of the Boltzmann equations shows that oscillations rather start when $(\delta m)^2/H = \Gamma_{scat}$, see Ref. [21]. In our scenario, as we also have checked from a numerical integration of the relevant Boltzmann equations, this turns out to happen at a temperature above T_{EW} . Thus we conclude that oscillations start as soon as EWSB occurs. As a result, annihilations can restart from this temperature and to determine how much of them will annihilate, one can just take the Boltzmann equations with the oscillation and annihilation terms,

$$\frac{d\Sigma_{\eta^0}}{dz} = -\frac{\langle\sigma_0 v\rangle s}{2zH} \left[\Sigma_{\eta^0}^2 - \Delta_{\eta^0}^2 - \Xi_{\eta^0}^2 - \Sigma_{\eta^0}^{eq^2} \right], \quad (24)$$

$$\frac{d\Delta_{\eta^0}}{dz} = 2i \frac{\delta m}{zH} \Xi_{\eta^0}, \quad (25)$$

$$\frac{d\Xi_{\eta^0}}{dz} = 2i \frac{\delta m}{zH} \Delta_{\eta^0} - \frac{\langle\sigma_0 v\rangle s}{zH} \Xi_{\eta^0} \Sigma_0, \quad (26)$$

where for any $T \leq T_{EW}$ we define $z \equiv m_{H^0}/T$, with $\langle\sigma_0 v\rangle$ the thermally averaged $\eta^0 \eta^{0*} \rightarrow SM$ annihilation cross section, and Ξ_{η^0} a quantity that accounts for the coherence between the η^0 and η^{0*} components (see [21] for further details). The resolution of these equations leads to a monotonically decreasing $\Sigma_{\eta^0}(z)$ function and to oscillating functions $\Delta_{\eta^0}(z) \propto \cos[f(z)]$ and $\Xi_{\eta^0}(z) \propto \sin[f'(z)]$ whose amplitudes also decrease monotonically. For fast oscillations, and neglecting the $\Sigma_{\eta^0}^{eq}$ term in Eq. (24), the set of Boltzmann equations can be simplified and solved analytically, at an approximate level, as explained in the Appendix. The solution it gives for Σ_{η^0} is ⁴

$$\Sigma_{\eta^0}(z \geq z_{EW}) = \frac{\Sigma_{\eta^0}(z_{EW})}{1 + \frac{1}{2} \frac{\langle\sigma_0 v\rangle s(z)}{H(z)} \left(\frac{z}{z_{EW}} - 1 \right) \Sigma_{\eta^0}(z_{EW})}, \quad (27)$$

with $z_{EW} = m_{H^0}/T_{EW}$ and where we fixed the initial abundance and asymmetry to be equal to $\Sigma_{\eta^0}(z_{EW}) = \Delta_{H^0}^{\lambda_s}/2$. The asymmetry Δ_{η^0} and Ξ_{η^0} are, in turn, fast oscillatory functions which are equal to zero on average.

⁴ This result is approximately the same than the one obtained in [21] for much smaller δm values – see Eqs (25) and (33) therein – but in which $x_{osc,ann}$ (which depends on δm) is now simply replaced by z_{EW} .

The result of Eq. (27) can also be qualitatively understood in the following way. Once $T \leq T_{EW}$, the fast oscillations reprocess quasi instantaneously the η^0 asymmetry in oscillatory abundances for η^0 and η^{0*} . On average, just after EWSB, we have therefore $n_{\eta^0} \simeq n_{\eta^{0*}} \simeq |\Delta n_{\eta^0}|_{T_{EW}}/2$. Since $T_{fo} > T_{EW}$, when two conjugate particles annihilate to SM particles, the reduction of inert doublet state it implies will not be compensated by any inverse processes. As a result, the Boltzmann equation for Σ_{η^0} one gets along this way is simply given by

$$\frac{d\Sigma_{\eta^0}}{dz} = -\frac{\langle\sigma_0 v\rangle s}{2zH} \Sigma_{\eta^0}^2, \quad (28)$$

whose resolution leads to nothing else than Eq. (27).⁵

The next step is to include the contribution of the charged states. Since these states do not oscillate, one could naively expect that the charged asymmetry is essentially left intact until the charged states decay to H^0 states. This doesn't work this way. To see that precisely, one should in principle solve the corresponding set of six coupled Boltzmann equations, for Σ_{η_0} , Δ_{η_0} , Ξ_{η_0} , Π_{η_0} , Σ_{η^+} , Δ_{η^+} , where Π_{η_0} and Ξ_{η_0} are the real and imaginary parts of the quantity that accounts for the coherence effects. Nevertheless in practice we don't need to go that far. It turns out that, if just before EWSB there are essentially only η^- and η^{0*} states as considered here, as soon as oscillations start they put the neutral state asymmetry to zero (on average), and processes which can transfer a charged asymmetry into a neutral one will very quickly put the charged asymmetry to zero too. This will be done in particular by $\eta^{0(*)} SM \leftrightarrow \eta^\pm SM$ inelastic scatterings and $\eta^\pm \leftrightarrow \eta^{0(*)} SMSM$ decays. The decrease of the charged component asymmetry due to these processes is exponential ($\Delta_{\eta^+} \propto e^{-z}$, as can be seen from the corresponding term in the Boltzmann equation, $sHz d\Delta_{\eta^+}/dz \propto -\Delta_{\eta^+} \gamma_{\eta^+ \leftrightarrow \eta^0} + \dots$). This “re-equilibration” of the asymmetries by these processes, which follows their “desilagnement” by the oscillations when these latter start, occurs much faster than the process of suppression of Σ_{η_0} in Eq. (50). As a result, in the same way as for the neutral states, one can adopt the simple assumption that as soon as oscillations start, the particle and antiparticle densities for charged states are equilibrated, $Y_{\eta^0} = Y_{\eta^{0*}} = Y_{\eta^+} = Y_{\eta^-}$. At this point, the annihilation processes such as $\eta^+ \eta^- \rightarrow SMSM$, $\eta^+ \eta^{0*} \rightarrow SMSM$ and $\eta^- \eta^0 \rightarrow SMSM$ can start again, in the same way as the $\eta^0 \eta^{0*} \rightarrow SMSM$ ones. The whole effect can be approximatively accounted by the simple Boltzmann equa-

⁵ The reason why the two results coincide is in fact more subtle. Since the $\eta^{0(*)}$ oscillatory behavior is given by

$$Y_{\eta^{0(*)}} = \frac{1}{2} f(z) (1 \pm \cos g(z)),$$

the Boltzmann equation in this naive approach should read

$$\frac{d\Sigma_{\eta^0}}{dz} = -2 \frac{\langle\sigma_0 v\rangle s}{zH} Y_{\eta^0} Y_{\eta^{0*}} = -\frac{\langle\sigma_0 v\rangle s}{2zH} \Sigma_{\eta^0}^2 \sin^2 g(z).$$

Averaging this expression, we find Eq. (28) up to an extra factor 1/2. An extra factor 2 must nevertheless be added to take into account the contribution of the coherence Ξ_{η^0} part, giving back Eq. (27).

tion

$$\frac{d\Sigma_{H_2}}{dz} = -\frac{\langle\sigma_{eff}v\rangle s}{zH} \Sigma_{H_2}^2. \quad (29)$$

Similarly to what has been obtained in Eq. (27), the resolution of Eq. (29), integrated from T_{EW} until now and using the initial condition in (18), leads to

$$\Sigma_{H_2}(z \geq z_{EW}) = \frac{\Delta_{H_2}^{\lambda_5}}{1 + \frac{\langle\sigma_{eff}v\rangle s(z)}{H(z)} \left(\frac{z}{z_{EW}} - 1\right) \Delta_{H_2}^{\lambda_5}}. \quad (30)$$

This equation holds for the case where the total number density just before EWSB is given by the asymmetry. If there is also a non-negligible part which is left from the symmetric freeze-out, one must simply replace the asymmetry at z_{EW} , $\Delta_{H_2}^{\lambda_5}$, by the total number density at the same temperature, $\Sigma_{H_2}(z_{EW})$, since this is the number which determines the number of symmetric annihilations which will occur after EWSB,⁶

$$\Sigma_{H_2}(z \geq z_{EW}) = \frac{\Sigma_{H_2}(z_{EW})}{1 + \frac{\langle\sigma_{eff}v\rangle s(z)}{H(z)} \left(\frac{z}{z_{EW}} - 1\right) \Sigma_{H_2}(z_{EW})}. \quad (31)$$

with $\Sigma_{H_2}(z_{EW}) \simeq \Sigma_{H_2}(z_{fo})$ as given in Eq. (16).

Note also that the $\eta^0 SM \leftrightarrow \eta^+ SM$ (and conjugated) processes above not only equilibrate the neutral and charged asymmetries, but also can break the coherence of the η^0 - η^{0*} by transforming a coherent neutral state into a charged state which does not oscillate. However, in the same way as for the Z exchange channel above, its rate goes under $(\delta m)^2/H$ before EWSB occurs, so that they do not prevent oscillations to start at $T = T_{EW}$.

Finally, because of the mass splittings between H^0 and the other components of the inert doublet, this total density is progressively transferred into a H^0 density through the decays of the heavier components. Note that, as we will see, in the numerical section below, Eq. (31) reaches its asymptotic value to a good approximation before T drops below the value of the mass splitting $m_{A^0} - m_{H^0}$. As a result, this splitting can be neglected as it was done to get Eq. (31).

F. Final inert scalar relic density

We summarize in Fig. 1 the evolution of the asymmetry $|\Delta_{H_2}|$ and total density Σ_{H_2} . We remind the main steps:

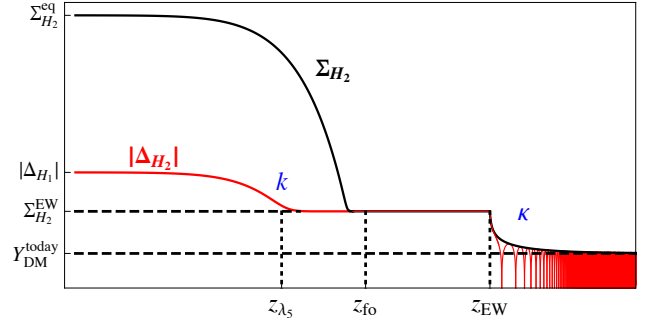


FIG. 1: Sketch of the scenario considered in section II and III. We represent, as a function of $z = m_{H^0}/T$, the H_2 asymmetry $|\Delta_{H_2}|$ in red and the total DM density Σ_{H_2} in black. We neglect in this sketch the mass splittings between the different components of H_2 . First step: The initial asymmetry Δ_{H_1} and Δ_{H_2} are fixed by the $B-L$ asymmetry. Second step: this asymmetry gets suppressed until the λ_5 interactions decouple at z_{λ_5} — the suppression is characterized by the k -factor. Third step: at z_{λ_5} the total annihilation cross section is still in thermal equilibrium and $\Sigma_{H_2}^{eq}$ follows the equilibrium density $\Sigma_{H_2}^{eq}$ until it reaches $|\Delta_{H_2}|$. Fourth step: at EWSB, oscillations start and reequilibrate the particle-antiparticle populations. At this point, annihilations can start again and if they do they deplete the density — the suppression is characterized by the κ -factor.

1. $T \gtrsim m_{H_2}$. The H_2 asymmetry, proportional to the $B-L$ asymmetry, is generated through the λ_5 interactions:

$$\Delta_{H_2}(z \lesssim 1) = -\frac{4}{23} \Delta_{B-L}. \quad (7)$$

2. $m_{H_2} \gtrsim T \geq T_{\lambda_5}$. The asymmetry undergoes a Boltzmann suppression until the λ_5 interaction decouples,

$$\Delta_{H_2}^{\lambda_5} \equiv |\Delta_{H_2}(z_{\lambda_5})| = \frac{16k(z_{\lambda_5})}{158 + 13k(z_{\lambda_5})} \Delta_{B-L}. \quad (12)$$

3. $T_{\lambda_5} > T > T_{EW}$. The symmetric component of Σ_{H_2} follows an exponential suppression, until Σ_{H_2} reaches at $z = z_{fo}$ the value

$$\Sigma_{H_2}(z_{fo}) \simeq \left[\Delta_{H_2}^{\lambda_5 2} + \Sigma_{H_2}^{eq 2}(z_{fo}) \right]^{1/2}. \quad (16)$$

From z_{fo} to z_{EW} , the annihilations are momentarily frozen, so that $\Sigma_{H_2}(z_{EW}) \simeq \Sigma_{H_2}(z_{fo})$. During this period, if the contribution from usual freeze-out is negligible, there are only \bar{H}_2 particles in the plasma and $\Sigma_{H_2}(z_{EW}) \simeq \Delta_{H_2}^{\lambda_5}$.

4. $T < T_{EW}$. The fast $\eta^0 \leftrightarrow \eta^{0*}$ oscillations start. They quickly reprocess the \bar{H}_2 density in equal abundances (on average) for η^0 , η^{0*} , η^+ and η^- . The annihilations can therefore start again and, if they do, they deplete the set of densities, whose sum reads asymptotically

$$\Sigma_{H_2}(z \gg z_{EW}) = \frac{\Sigma_{H_2}(z_{EW})}{1 + \frac{\langle\sigma_{eff}v\rangle s(z)}{H(z)} \frac{z}{z_{EW}} \Sigma_{H_2}(z_{EW})}. \quad (32)$$

⁶ If there is no asymmetry and if the freeze-out has occurred prior to EWSB, one recognizes in Eq. (31) the usual asymptotic freeze-out behavior, i.e. the freeze-out is not instantaneous, but reaches asymptotically its final value as given in this equation. In practice, as well known, the effect is negligible in this case, i.e. the denominator is equal to unity to a good approximation. Here, instead, the denominator at z_{fo} can be much larger due to the asymmetry.

This total density is progressively transferred into a H^0 density through the decay of the heavier components.

The final DM abundance is therefore given by

$$Y_{DM}^{today} = \Sigma_{H_2}(z \gg z_{EW}) = \frac{\Sigma_{H_2}(z_{EW})}{1 + \kappa \cdot \Sigma_{H_2}(z_{EW})}, \quad (33)$$

where we define

$$\kappa \equiv \frac{\langle \sigma_{effv} \rangle s(z)}{H(z)} \frac{z}{z_{EW}} \simeq 1.3 \cdot 10^{13} \cdot \left(\frac{\langle \sigma_{effv} \rangle}{1 \text{ pb}} \right). \quad (34)$$

Since ultimately no asymmetry survives, the relation between the baryon and the $B - L$ asymmetry is still given by Eq. (21), and the final DM to baryon density ratio is given by

$$\frac{\Omega_{DM}}{\Omega_B} = \frac{\Sigma_{H_2}(z_{EW})}{1 + \kappa \cdot \Sigma_{H_2}(z_{EW})} \cdot \frac{1}{Y_B^{today}} \cdot \left(\frac{m_{H^0}}{1 \text{ GeV}} \right), \quad (35)$$

or equivalently, if the asymmetric component dominates, using Eqs. (12) and (18),

$$\frac{\Omega_{DM}}{\Omega_B} = \frac{148k(z_{\lambda_5})}{474 + (39 + 148\kappa Y_B^{today})k(z_{\lambda_5})} \cdot \left(\frac{m_{H^0}}{1 \text{ GeV}} \right). \quad (36)$$

A number of comments can be done regarding these results:

- Eqs. (33) and (35) show that beside the λ_5 interaction induced “ k -factor” suppression in Δ_{H_2} , see Eq. (12), oscillations drive a $1 + \kappa \cdot \Sigma_{H_2}(z_{EW})$ factor suppression. This “ κ -factor” suppression can be sizable as soon as $\kappa \cdot \Sigma_{H_2}(z_{EW}) \gtrsim 1$.
- As Eq. (31) shows, this suppression is neither instantaneous nor exponential. It goes as the inverse of $z/z_{EW} - 1$ until it reaches an asymptotic value. In this sense, imposing that the cross section satisfies the unitarity bound, it is naturally limited but it still can be responsible for the $\sim (10 \text{ GeV}/m_{DM})$ suppression needed, see below.
- The appearance of the $\kappa \cdot \Sigma_{H_2}(z_{EW})$ factor is not surprising. The condition $\kappa \cdot \Sigma_{H_2}(z_{EW}) < 1$ is nothing but the condition $(n_{H_2} + n_{\bar{H}_2})\langle \sigma_{effv} \rangle < H$ at $T = T_{EW}$.
- Interestingly, for large values of $\kappa \cdot \Sigma_{H_2}(z_{EW})$, the Y_{DM}^{today} relic density obtained doesn’t depend anymore on the asymmetry left at T_{EW} , even if this asymmetry is the source of the final DM abundance. In this case, we simply get

$$Y_{DM}^{today} = \frac{1}{\kappa},$$

and the ratio reads

$$\frac{\Omega_{DM}}{\Omega_B} \simeq 0.15 \cdot z_{EW} \cdot \left(\frac{1 \text{ pb}}{\langle \sigma_{effv} \rangle} \right). \quad (37)$$

This means, as we could have anticipated, that for large cross section the asymmetry left is independent of the initial asymmetry, provided this initial asymmetry is large enough. In other terms, if the κ -factor suppression is small, both baryon and DM asymmetries are directly connected. If instead it is large, they are not related anymore in a so direct way, since in this case the final relic density depends only on the annihilation cross section.⁷ Note interestingly that Eq. (37) is nothing but the result of the standard freeze-out scenario, but with the important difference that in the standard case, z_{EW} in Eq. (37) must be replaced by z_{fo} .

III. FAILURE OF THE ASYMMETRIC IDM SCENARIO

The final result of Eq. (35) depends on three parameters: m_{H^0} , $\Sigma_{H_2}(z_{EW})$ and the total cross section $\langle \sigma_{effv} \rangle$ via κ in Eq. (34). This means that for given values of the input parameters m_{H^0} and $\langle \sigma_{effv} \rangle$, there is only one value of $\Sigma_{H_2}(z_{EW})$ which gives the observed value of Ω_{DM}/Ω_B , as given by the PLANCK best fits, $\Omega_{DM}h^2 = 0.120$ and $\Omega_B h^2 = 0.022$ [1]. Since $\Sigma_{H_2}(z_{EW})$ depends only on these two input parameters and on $\Delta_{H_2}^{\lambda_5}$, this means also that there is only one value of $\Delta_{H_2}^{\lambda_5}$ which gives the correct relic density for fixed values of the two input parameters. We show in Fig. 2 this value of $\Delta_{H_2}^{\lambda_5}$ as a function of m_{H^0} for different values of the cross section. By comparing this value of $\Delta_{H_2}^{\lambda_5}$ to the value this asymmetry would have if there were no “ k -factor” suppression – given by the Δ_{H_1} upper horizontal line – one can read off what is the value of this λ_5 induced “ k -factor” suppression, Eq. (12) as compared to Eq. (7).

As said above, to dominate the final relic density, the asymmetry cannot be suppressed by more than a factor $m_{DM}/10 \text{ TeV}$. Figure 2 also shows the corresponding values of the $\kappa \cdot \Sigma_{H_2}(z_{EW})$ factor which lead to the other suppression, i.e. the $1/(1 + \kappa \cdot \Sigma_{H_2}(z_{EW}))$ factor in Eq. (35). It also shows for which values of the various parameters the asymmetry produced before the EW transition is responsible for 50% of the final DM relic density (black line). Above (below) this line the relic density is dominantly of asymmetric (symmetric) origin. Similarly, the dotted upper (lower) black line gives the values of the parameters above (below) which the asymmetry is responsible for more (less) than 90% (10%) of the final relic density. For masses which give a freeze-out below T_{EW} , the $\kappa \cdot \Sigma_{H_2}(z_{EW})$ factor becomes exponentially large because in this case $\Sigma_{H_2}(z_{EW})$ is still exponentially larger than its value at freeze-out. Thus, the proportion of $\Sigma_{H_2}(z_{EW})$ which is due to $\Delta_{H_2}^{\lambda_5}$ is therefore exponentially suppressed. This explains why the black lines quickly go up for m_{DM} below 4 – 5 TeV. Note nevertheless that this suppression, even if exponential, is far from instantaneous. As a result we find that, still, the

⁷ But still, even in this case, they remain similar as the κ factor is bounded from above by unitarity considerations on the total cross section.

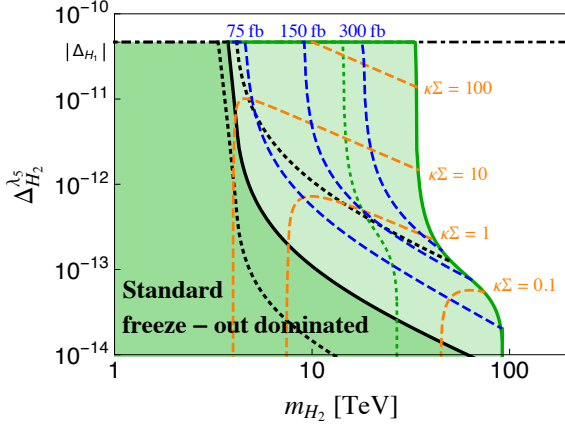


FIG. 4: Same as Fig. 2, but allowing the inert scalar to decay into a lighter real scalar singlet S with mass $m_S = m_{H_0}/10$.

maintaining asymmetries. Smaller values of λ_5 rather quickly lead to no thermalization of the H_2 and H_1 asymmetries, i.e. to no creation of a H_2 asymmetry. In most of the relic density allowed parameter space, both the “ k ” and “ κ ” suppressions are active, although it is possible to have only one of the effect to account for all the necessary suppression. As said above, an important constraint that one must satisfy is the direct detection constraint of Eq. (4). The value of the mass splitting just quoted are below the ~ 100 keV direct detection lower bound of Eq. (4).

Thus, unless direct detection would allow a mass splitting as low as the value $m_{A^0} - m_{H^0} \sim 15$ keV, which seems very unlikely, this very minimal asymmetric scenario is in fact excluded! This can also be clearly seen from Fig. (3) where the region allowed by direct detection taking in Eq. (4) a mass splitting $\delta m_{\min} = 100$ keV has no overlap with the region which gives the observed relic density. Or, in other words, imposing that the mass splitting is above 100 keV, the λ_5 interaction turns out to decouple only at $z_{\lambda_5} \gtrsim 50$ leading to a tiny Ω_{DM} relic density.

IV. REPROCESSING THE INERT DOUBLET ASYMMETRY INTO A LIGHTER PARTICLE DM RELIC DENSITY

Since the very minimal IDM scenario above cannot account for both the relic density and direct detection constraints at the same time, one question one must ask is whether this simple scenario of an IDM asymmetry creation could not be nevertheless at the origin of the DM relic density today in a simple way. This could in fact happen if the DM is made of a lighter specie, whose relic density would be due to the reprocessing of the inert doublet asymmetry into this specie. Such a reprocessing could for instance take place through decay. For the scalar scenario we consider, this could be the case if there exists a lighter Z_2 odd particle, “ S ”, to which the inert doublet states can decay. In this case, if the asymmetry is fully reprocessed into this lighter particle S , so that each inert scalar

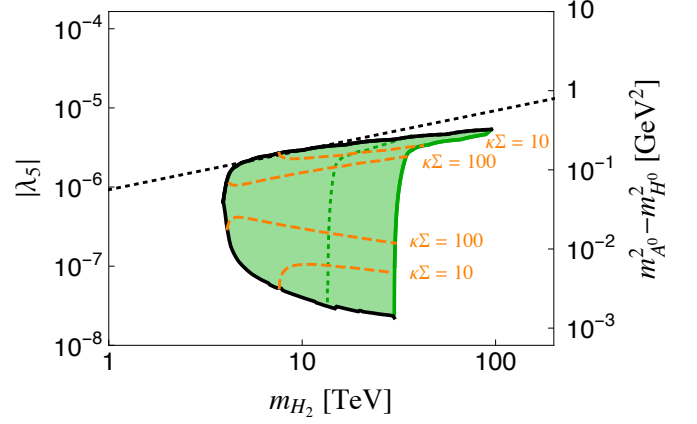


FIG. 5: Same as Fig. 3, but allowing the inert scalar to decay into a lighter real scalar singlet S with mass $m_S = m_{H_0}/10$.

component gives one S particle, the results of Figs. 2 and 3 are still fully valid provided the mass of the S particle, m_S , is close to m_{H_2} . If instead it is sizably smaller, this requires to create more inert particles by a factor m_{H_2}/m_S . As an example, Figs. 4 and 5 show the value of parameters we need to get the observed relic density for a ratio m_{H_2}/m_S equal to 10. Note that such large H_2 asymmetries cases, beside allowing smaller DM masses, also give relaxed lower bounds on the $\lambda_{3,4}$ couplings (in order to suppress sufficiently the symmetric part). Sizably smaller values of these couplings are possible, relaxing accordingly the Landau pole constraints. In order to reprocess the inert doublet asymmetry into such a S specie, various possibilities could be considered.

A simple possibility is to consider S as a Z_2 -odd scalar singlet, into which the scalar doublet states decay sufficiently slowly to happen after the freeze-out of this singlet DM particle. Such a decay can be accounted for by a $\mathcal{L} \ni \mu_S H_1^\dagger H_2 S$ renormalizable interaction. If so, the main constraint to satisfy along such a scenario is, in order that the S relic density is mainly produced from the IDM asymmetry, that the S particles has a $S^\dagger S \rightarrow SM SM$ annihilation channel with a large enough cross section to leave a relic density smaller than the observed one at S freeze-out. These annihilations can be accounted for by a $\mathcal{L} \ni \lambda_S H_1^\dagger H_1 S^\dagger S$ interaction.

As an example, if we take a real scalar singlet, and fix the parameters to be $m_S \sim 2$ TeV (400 GeV) and $m_{H_2} = 10$ TeV, both conditions are fulfilled for $\lambda_S \gtrsim 0.6(0.1)$ and $\mu_S \lesssim 4 \cdot 10^{-5}$ GeV ($7 \cdot 10^{-6}$ GeV). Also, the λ_S interaction induces elastic scattering on nucleon through SM scalar exchange, $\sigma_N = \lambda_S^2 m_N^4 f_N^2 / (\pi m_h^4 m_S^2)$, where m_N is the nucleon mass and the nucleon form factor is approximately given by $f_N \approx 0.3$. The LUX experiment constraint [28], which for $m_{DM} \gtrsim 100$ GeV is $\sigma_N \lesssim 1.2 \cdot 10^{-11} (m_S/1 \text{ GeV})$ pb, is satisfied if $\lambda_S \lesssim 1.6 \cdot 10^{-5} (m_S/1 \text{ GeV})^{3/2}$. Combining both these lower and upper bounds on λ_S leads to the lower bound $m_S \gtrsim 300$ GeV [29]. In Fig. 6 we show the evolution of the asymmetries we get as a function of the temperature for an example of parameter set which leads to the observed relic density.

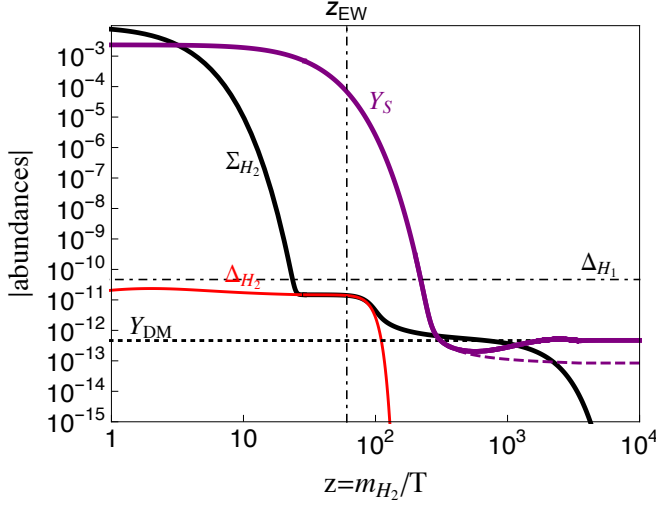


FIG. 6: Evolution of the various abundances as a function of $z = m_{H_2}/T$ in the case where the inert doublet decays into a real state S after S freezeout. Such an evolution has been obtained fixing the H_2 -related parameters to $m_{H_2} = 10$ TeV, $\langle\sigma_{eff}\rangle = 0.5$ pb and $\delta m_{H_2} = 3 \cdot 10^{-6}$ GeV (corresponding to $\lambda_5 \approx 10^{-6}$), the S -related parameters to $m_S = 1$ TeV and $\langle\sigma_{Sv}\rangle = 4$ pb, and the connector parameter controlling the decay rate to $\mu_S = 5 \cdot 10^{-6}$ GeV. The dashed purple curve shows what would be the evolution of the S density if there were no $H_2 \rightarrow H_1 S$ decay.

Similarly to the fermion scenario considered in Ref. [16], another possibility is to consider instead that the decays occur when the freeze-out of the singlet particle S has still not taken place. In this case, the inert doublet asymmetry could also be at the origin of the DM relic density, if the singlet is a complex field and if the inert doublet asymmetry is reprocessed into a S asymmetry. Since inert doublet oscillations start at T_{EW} , this requires the reprocessing to be done prior to EWSB. Imposing in addition for simplicity that the decay occurs after the λ_5 interaction has decoupled at z_{λ_5} , for example for $m_{H_2} = 10$ TeV and $m_S \sim 2$ TeV, one needs 10^{-5} GeV $\lesssim \mu_S \lesssim 10^{-3}$ GeV. For this scenario to work, one has to make sure that the S asymmetry created in this way is not washed-out by possible S - S^\dagger oscillations. This requires that terms as $\lambda'_5 H_1^\dagger H_1 (S^2 + h.c.)$ or $m_S^2 S^2 + h.c.$ are sufficiently suppressed for the oscillations not to occur before S freeze-out. This means the S mass splitting $\delta m_S = (m_S'^2 + \theta(T_{EW} - T)\lambda'_5 v^2)/2m_S$ must be smaller than

$$\delta m_S \lesssim 10^{-2} \cdot (z_{fo}^S)^{-5/2} \cdot \left(\frac{m_S}{1 \text{ TeV}}\right)^2 \cdot \sqrt{\frac{\langle\sigma_{Sv}\rangle}{1 \text{ pb}}} \text{ eV}, \quad (38)$$

with $z_{fo}^S \gtrsim 20$ the value of m_S/T at which the S freeze-out occurs, and $\langle\sigma_{Sv}\rangle$ the S annihilation cross section. Note that at temperature lower than T_{fo}^S , when the S oscillations starts at z_{osc}^S , they can allow the S annihilation to restart in the same way as for the inert doublet above. Similarly to Eq. (33), this causes a suppression of the S asymmetry by a factor equal to $(1 + \kappa_S \Sigma_S(z_{fo}^S))^{-1}$ with $\kappa_S = \langle\sigma_{Sv}\rangle sz/H(z)z_{osc}^S$. In Fig. 7 we show an example of evolution of the H_2 and S asymmetries along such a scenario.

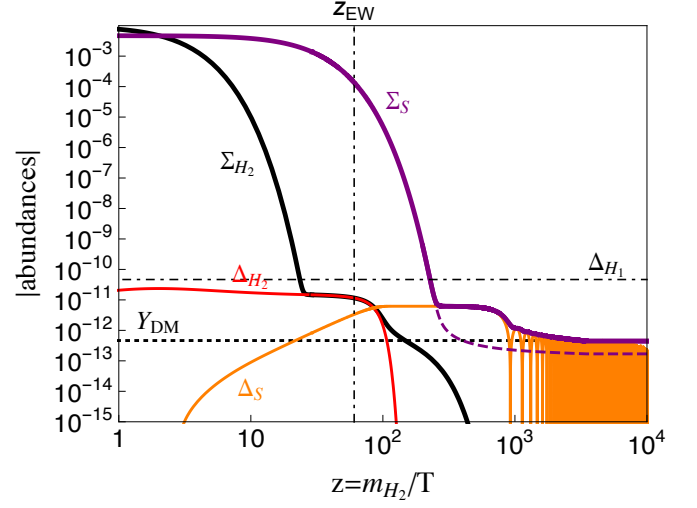


FIG. 7: Evolution of the various abundances as a function of $z = m_{H_2}/T$ in the case where the inert doublet decays into a complex state S before S freezeout, for $m_{H_2} = 10$ TeV, $\langle\sigma_{eff}\rangle = 1$ pb, $\delta m_{H_2} = 3 \cdot 10^{-6}$ GeV (corresponding to $\lambda_5 = 10^{-6}$), $m_S = 1$ TeV, $\langle\sigma_{Sv}\rangle = 4$ pb, $\delta m_S = 10^{-7}$ eV, and $\mu_S = 5 \cdot 10^{-5}$ GeV. The dashed purple curve shows what would be the evolution of the S density if there were no $H_2 \rightarrow H_1 S$ decay.

V. SUMMARY

In summary, if there exists an inert scalar doublet H_2 , unless the λ_5 interaction is tiny, the inert doublet components will automatically develop an asymmetry from thermalization with the ordinary SM scalar doublet and lepton asymmetries. We have studied in details what is the fate of such an asymmetry at temperature below the value of the inert doublet scalar mass, m_{H_2} . Beside being responsible for the asymmetry creation, the λ_5 interaction also controls the neutral component mass splitting (hence the Z -exchange direct detection rate) and induces a “ k -factor” suppression of the inert doublet asymmetry at temperature below m_{H_2} . On top of this suppression one can also have an extra “ κ -factor” suppression, from the combined effect of DM oscillation (also induced by the λ_5 interaction) and DM symmetric annihilation. This leads to a scenario which chronologically occurs as represented in Fig. 1. We showed that in the few-TeV range, there is an all region of parameter space where the DM asymmetry survives enough to lead to the observed DM relic density, Fig. 2, but this region turns out to lead to a too large Z -exchange direct detection contribution. As a result this scenario is nothing but excluded.

Next we looked at the possibility that the inert scalar asymmetry produced could still be at the origin of the observed DM relic density, which could be the case if it is reprocessed to a lighter specie, S , which satisfies the direct detection constraints. We considered 2 scenarios where DM is made of a singlet odd under the Z_2 symmetry. a) Slow decay of the asymmetry into the (real or complex) singlet particle, occurring after S freeze-out. b) Reprocessing of the inert scalar doublet asymmetry into a S asymmetry through faster decays occurring before S freeze-out. Both possibilities can lead to

the observed relic density provided the interaction causing the decay is small enough to induce this decay at the right time.

As most asymmetric DM scenarios, the framework we consider does not explain why the baryon and DM abundances are so similar. Our scenario trades this abundance coincidence for a coincidence between the mass of the proton, the mass of the inert states, the mass of the dark matter particle, and the values of various couplings. Even if both abundances have same origin, these parameters must "cooperate" to lead to a DM abundance so close to the baryon one. Rather than providing a real explanation for the abundance coincidence, this scenario shows instead that the origin of the DM relic density could be of asymmetric origin, due to the generation of an inert scalar asymmetry related to the generation of a $B - L$ asymmetry at high temperature.

Acknowledgement

We acknowledge useful discussions with M. Cirelli, G. Servant, S. Tulin and M. Tytgat. This work is supported by the FNRS-FRS, the IISN, the FRiA and an ULB-ARC and the Belgian Science Policy, IAP VI-11.

Appendix

Rates and cross sections

In Eq. (8), the reaction density of the λ_5 scatterings for the η^+ (and similarly for η^0) is given by

$$\gamma_{\lambda_5} = \gamma_{\eta\eta}^{\phi\phi} + \gamma_{\phi\eta}^{\phi\eta}, \quad (39)$$

where

$$\gamma_{cd}^{ab} = \frac{m_{H_2}^4}{64\pi z} \int_4^\infty dx \sqrt{x} K_1(z\sqrt{x}) \hat{\sigma}(ab \rightarrow cd). \quad (40)$$

with $\hat{\sigma}(ab \rightarrow cd)$ the reduced cross section. These are given by

$$\hat{\sigma}_{\lambda_5}^s(\phi\phi \rightarrow \eta\eta) = \frac{3\lambda_5^2}{2\pi} \sqrt{1 - \frac{4}{x}}, \quad (41)$$

$$\hat{\sigma}_{\lambda_5}^t(\phi\eta \rightarrow \phi\eta) = \frac{3\lambda_5^2}{2\pi} \left(1 - \frac{1}{x}\right)^2. \quad (42)$$

In the non-relativistic limit, the corresponding rate is given by

$$\Gamma_{\lambda_5} \equiv \frac{\gamma_{\lambda_5}}{n_{\eta^+}^{eq}} \equiv n_{\eta^+}^{eq} \langle \sigma_{\lambda_5}^s v \rangle + n_{\phi^+}^{eq} \langle \sigma_{\lambda_5}^t v \rangle, \quad (43)$$

where

$$\langle \sigma_{\lambda_5}^s v \rangle = \frac{3\lambda_5^2}{32\pi m_{H_2}^2}, \text{ and } \langle \sigma_{\lambda_5}^t v \rangle = \frac{3\lambda_5^2}{16\pi m_{H_2}^2}. \quad (44)$$

In Eq. (15) and (29), the effective cross section of the $H_2 \bar{H}_2 \rightarrow SM SM$ coannihilations is given by [9]

$$\langle \sigma_{eff} v \rangle = \sum_{i,j} \langle \sigma_{ij} v \rangle \frac{Y_i^{eq} Y_j^{eq}}{\Sigma_{H_2}^{eq} \Sigma_{H_2}^{eq}} \simeq \frac{1}{64\pi m_{H_2}^2} \left(\frac{3}{8} g^4 + \lambda_3^2 + \lambda_4^2 \right). \quad (45)$$

where g is the weak coupling constant. We neglected the λ_5 contribution, and the corrections due to the contributions proportional to $\langle v^2 \rangle$.

Analytical resolution of the Boltzmann equations

The Boltzmann equations given in Eqs. (24)-(26) do not in general have a simple analytical solution. However, in the case of very fast oscillations, like it is the case here, a good approximation consists in symmetrizing the equations for Δ_{η^0} and Ξ_{η^0} , i.e. replacing Eqs. (25)-(26) by

$$\frac{d\Delta_{\eta^0}}{dz} = 2i \frac{\delta m}{zH} \Xi_{\eta^0} - \frac{1}{2} \frac{\langle \sigma_0 v \rangle s}{zH} \Delta_{\eta^0} \Sigma_{\eta^0}, \quad (46)$$

$$\frac{d\Xi_{\eta^0}}{dz} = 2i \frac{\delta m}{zH} \Delta_{\eta^0} - \frac{1}{2} \frac{\langle \sigma_0 v \rangle s}{zH} \Xi_{\eta^0} \Sigma_{\eta^0}. \quad (47)$$

In this approximation, the solutions for Δ_{η^0} and Ξ_{η^0} are of the form

$$\Delta_{\eta^0}(z) = f(z) \cos[g(z)], \quad \Xi_{\eta^0}(z) = i f(z) \sin[g(z)]. \quad (48)$$

Furthermore, since we are interested in oscillations happening after the freeze-out, we can neglect $\Sigma_{\eta^0}^{eq}$ in Eq. (24). With these approximations, integrating from z_{EW} to z with the initial conditions $\Delta_{\eta^0}(z_{EW}) = \Sigma_{\eta^0}(z_{EW})$ and $\Xi(z_{EW}) = 0$, the analytical solutions of the Boltzmann equations Eqs. (24)-(47) are given by Eq. (48) and

$$\Sigma_{\eta^0}(z) = \sqrt{\Delta_{\eta^0}^2(z) - \Xi_{\eta^0}^2(z)} = f(z), \quad (49)$$

with

$$f(z) = \frac{\Sigma_{\eta^0}(z_{EW})}{1 + \frac{1}{2} \frac{\langle \sigma_0 v \rangle s(z)}{H(z)} \left(\frac{z}{z_{EW}} - 1 \right) \Sigma_{\eta^0}(z_{EW})}, \quad (50)$$

$$g(z) = \frac{\delta m}{H(z)} \left(\frac{z^2}{z_{EW}^2} - 1 \right). \quad (51)$$

The abundance Σ_{η^0} decreases therefore monotonically until it reaches an asymptotical value given by

$$\Sigma_{\eta^0}(z \gg z_{EW}) = \frac{\Sigma_{\eta^0}(z_{EW})}{1 + \frac{1}{2} \frac{\langle \sigma_0 v \rangle s(z)}{H(z)} \frac{z}{z_{EW}} \Sigma_{\eta^0}(z_{EW})}. \quad (52)$$

Note that despite appearance, the denominator doesn't depend on z , since $sz/(H z_{EW}) = 12\sqrt{g_*} M_{Pl} T_{EW}/5\pi^2$.

-
- [1] P. A. R. Ade *et al.* [Planck Collaboration], arXiv:1502.01589 [astro-ph.CO].
 - [2] H. Davoudiasl and R. N. Mohapatra, New J. Phys. **14** (2012) 095011 [arXiv:1203.1247 [hep-ph]].
 - [3] K. Petraki and R. R. Volkas, Int. J. Mod. Phys. A **28** (2013) 1330028 [arXiv:1305.4939 [hep-ph]].
 - [4] K. M. Zurek, Phys. Rept. **537** (2014) 91 [arXiv:1308.0338 [hep-ph]].
 - [5] S. M. Boucenna and S. Morisi, Front. Phys. **1** (2014) 33 [arXiv:1310.1904 [hep-ph]].
 - [6] R. Barbieri, L. J. Hall and V. S. Rychkov, Phys. Rev. D **74** (2006) 015007 [hep-ph/0603188].
 - [7] L. Lopez Honorez, E. Nezri, J. F. Oliver and M. H. G. Tytgat, JCAP **0702** (2007) 028 [hep-ph/0612275].
 - [8] E. Ma, Phys. Rev. D **73** (2006) 077301 [hep-ph/0601225].
 - [9] T. Hambye, F.-S. Ling, L. Lopez Honorez and J. Rocher, JHEP **0907** (2009) 090 [Erratum-ibid. **1005** (2010) 066] [arXiv:0903.4010 [hep-ph]].
 - [10] N. Nagata and S. Shirai, arXiv:1411.0752 [hep-ph].
 - [11] M. Cirelli, N. Fornengo and A. Strumia, Nucl. Phys. B **753** (2006) 178 [hep-ph/0512090].
 - [12] N. Fonseca, R. Z. Funchal, A. Lessa and L. Lopez-Honorez, arXiv:1501.05957 [hep-ph].
 - [13] M. D’Onofrio, K. Rummukainen and A. Tranberg, Phys. Rev. Lett. **113** (2014) 14, 141602 [arXiv:1404.3565 [hep-ph]].
 - [14] J. A. Harvey and M. S. Turner, Phys. Rev. D **42** (1990) 3344.
 - [15] C. Arina and N. Sahu, Nucl. Phys. B **854** (2012) 666 [arXiv:1108.3967 [hep-ph]].
 - [16] G. Servant and S. Tulin, Phys. Rev. Lett. **111** (2013) 15, 151601 [arXiv:1304.3464 [hep-ph]].
 - [17] T. R. Dulaney, P. Fileviez Perez and M. B. Wise, Phys. Rev. D **83** (2011) 023520 [arXiv:1005.0617 [hep-ph]].
 - [18] S. M. Boucenna, M. B. Krauss and E. Nardi, arXiv:1503.01119 [hep-ph].
 - [19] S. Davidson, R. Gonzalez Felipe, H. Serdio and J. P. Silva, JHEP **1311** (2013) 100 [arXiv:1307.6218 [hep-ph]].
 - [20] E. Nardi, F. Sannino and A. Strumia, JCAP **0901** (2009) 043 [arXiv:0811.4153 [hep-ph]].
 - [21] M. Cirelli, P. Panci, G. Servant and G. Zaharijas, JCAP **1203** (2012) 015 [arXiv:1110.3809 [hep-ph]].
 - [22] S. Tulin, H. B. Yu and K. M. Zurek, JCAP **1205** (2012) 013 [arXiv:1202.0283 [hep-ph]].
 - [23] Y. Burnier, M. Laine and M. Shaposhnikov, JCAP **0602** (2006) 007 [hep-ph/0511246].
 - [24] M. Klasen, C. E. Yaguna and J. D. Ruiz-Alvarez, Phys. Rev. D **87** (2013) 075025 [arXiv:1302.1657 [hep-ph]].
 - [25] D. S. Akerib *et al.* [LUX Collaboration], Phys. Rev. Lett. **112** (2014) 9, 091303 [arXiv:1310.8214 [astro-ph.CO]].
 - [26] S. D. McDermott, H. B. Yu and K. M. Zurek, Phys. Rev. D **85** (2012) 023519 [arXiv:1103.5472 [hep-ph]].
 - [27] C. Kouvaris and P. Tinyakov, Phys. Rev. Lett. **107** (2011) 091301 [arXiv:1104.0382 [astro-ph.CO]].
 - [28] D. S. Akerib *et al.* [LUX Collaboration], Phys. Rev. Lett. **112** (2014) 091303 [arXiv:1310.8214 [astro-ph.CO]].
 - [29] J. M. Cline, K. Kainulainen, P. Scott and C. Weniger, Phys. Rev. D **88** (2013) 055025 [arXiv:1306.4710 [hep-ph]].